Low-energy photon–neutrino inelastic processes beyond the Standard Model

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Abstract

We investigate in this work the leading contributions of the MSSM with R-parity violation and of Left–Right models to the low-energy five-leg photon–neutrino processes. We discuss the results and compare them to the Standard Model ones.

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I. INTRODUCTION

The low-energy photon–neutrino interactions are of potential interest for astrophysics and cosmology (see [1–7] and references therein). By low energies, we mean that the photons and the neutrinos carry energies that do not exceed the confinement scale. The processes involving two neutrinos and two photons are strongly suppressed, not only because of the weak interaction but also because of Yang's theorem [8], which forbids a two photon-coupling to a $J = 1$ state.

The processes involving one more photon, such as

$$\gamma\nu \rightarrow \gamma\gamma\nu \tag{1}$$

$$\nu\bar{\nu} \rightarrow \gamma\gamma\gamma \tag{2}$$

$$\gamma\gamma \rightarrow \gamma\nu\bar{\nu}\,, \tag{3}$$

are not longer constrained by Yang's theorem. Moreover, the extra α in the five-leg cross section is compensated by the replacement of the ω/M_W suppression in the four-leg processes by an ω/m_e enhancement, ω being the centre-of-mass energy of the collision.

Recently, processes (1), (2) and (3) have been studied, first within an effective theory, see [4] and [6], based on the four-photon Euler–Heisenberg Lagrangian [9], which describes four-photon interactions. This effective approach gives reliable results for energies below the threshold of e^+e^- pair production, while the necessary energy, interesting for the study of the supernova dynamics, is above 1 MeV. The extrapolation to energies above 1 MeV being suspect, the processes cited above have been computed directly in the Standard Model, see [7] and [10].

For astrophysical implications, processes (1) and (2) may give contributions to the neutrino mean free path inside the supernova, and it is possible that process (3) is an energy loss mechanism for stars.

The conclusion we arrived at in [7] was that the five-leg photon–neutrino processes should be incorporated in supernova codes while they do not seem to have any relevance in cosmology. Indeed, we have found that these processes are unlikely to be important for the study of the neutrino decoupling temperature [7].

Knowing the contributions of these processes in the Standard Model, it is worth investigating whether they could have some importance and implications beyond the Standard Model.

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Among possible minimal supersymmetric (MSSM) extensions of the Standard Model (SM), the one including R-parity violating ($\not{R}_p$) processes has been attracting increasing attention over the last years [11]. In this paper, we compute the contribution of the process $\gamma\nu \to \gamma\gamma\nu$ in the MSSM, with $\not{R}_p$.

Another popular extension of the SM of the electroweak interactions is based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{\tilde{Y}}$ [12], in which we also compute the process $\gamma\nu \to \gamma\gamma\nu$.

It is worth pointing out that the range of energy in which we will apply the above reaction is well below the mass of the "exchanged particle", i.e. $W^{\pm}$ in the case of the Standard Model, $W^{\pm}$ or $W'^{\pm}$ in the Left–Right model (LR), and sleptons in the MSSM with $\not{R}_p$, so that the neutrino–electron coupling is treated as a four-Fermi interaction. In all cases, we perform the calculation with massless neutrinos.

At low energies, far from the confinement scale ~ 1 GeV, the leading contribution is given by diagrams involving only electrons or muons running inside the loop (see Fig. 1). It is precisely the appearance of m_e (or m_{μ}) as a scale, instead of M_W (which is the scale governing the four-leg photon–neutrino reactions in the SM), that makes such five-leg processes relevant at energies of the order of a few tens of MeV, in the SM. In the MSSM, gauge invariance forbids the process of Fig. 1a with a neutral scalar (sneutrino $\tilde{\nu}$) exchange; only Fig. 1b will thus contribute with a slepton $\tilde{l}$ exchange, as we will see in section II. In the Left–Right model, the bosons Z' and $W'^{\pm}$ enter the process in Figs. 1a and 1b.

The outline of the paper is as follows: section II is devoted to the computation of the five-leg process in the MSSM with $\not{R}_p$, while the computation reported in section III is done in the LR model. Finally, we end up with a discussion on the comparison of these contributions with the one obtained in the Standard Model.

II. COMPUTATION OF $\gamma\nu \to \gamma\gamma\nu$ IN THE MSSM WITH $\not{R}_P$

In SUSY extensions, gauge invariance and renormalizability no longer ensure lepton number L (or baryon number B) conservation. The generalization of the MSSM, which includes R-parity[1] violation ($\not{R}_p$), allows the interaction $\gamma\nu \to \gamma\gamma\nu$ to proceed via the exchange of a slepton, unlike the $W^{\pm}$ and Z boson exchange that takes place in the Standard

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[1]R_p is a multiplicative quantum number defined by $R_p = (-1)^{L+3B+2S}$, where L, B, S are the lepton number, baryon number and spin of the particle, respectively.

Model.

The relevant $\not{R}_p$ superpotential, consistent with Lorentz invariance, gauge and SUSY symmetries [11], is, in what concerns our *low-energy* process:

$$\mathcal{W} = \frac{1}{2}\lambda_{ijk}L^i L^j \bar{E}^k \ , \tag{4}$$

where i, j and k are family indices and L^i, E^i are the left-handed lepton and right-handed singlet charged lepton superfield, respectively. The antisymmetry under $SU(2)_L$ implies that the Yukawa couplings are also antisymmetric under the exchange of the first two family indices, namely $\lambda_{ijk} = -\lambda_{jik}$ and as far as our analysis is concerned, λ is assumed to be real. All the following bounds on the couplings λ_{ijk} are derived within the degenerate mass of sleptons and squarks [13] of $\tilde{m} \sim 100$ GeV:

$$\lambda_{12k} < 0.05 \left(\frac{\tilde{m}}{100 \text{ GeV}}\right) \quad \text{for} \quad k = 1, 2, 3 \ ,$$

$$\lambda_{13k} < 0.06 \left(\frac{\tilde{m}}{100 \text{ GeV}}\right) \quad \text{for} \quad k = 1, 2 \quad \text{and} \quad \lambda_{133} < 0.004 \left(\frac{\tilde{m}}{100 \text{ GeV}}\right),$$

$$\lambda_{23k} < 0.06 \left(\frac{\tilde{m}}{100 \text{ GeV}}\right) \quad \text{for} \quad k = 1, 2, 3 \ . \tag{5}$$

Expressing the relevant part of the superpotential in (4) in terms of component fields, we obtain the Lagrangian:

$$\mathcal{L}_{LLE} = \lambda_{ijk}[\ \tilde{\nu}_L^i \ e_L^j \ \bar{e}_R^k \ + \ \left(\tilde{e}_R^k\right)^c \ \nu_L^i \ e_L^j \ - \ \tilde{e}_L^i \ \nu_L^j \ \bar{e}_R^k \ ] \ + \text{h.c.} \tag{6}$$

Notice that there is no term containing a neutrino ν and a sneutrino $\tilde{\nu}$ field simultaneously. This implies that the process with a sneutrino exchanged in Fig. 1a is not allowed by gauge invariance and that only the last two terms in Eq. (6) will contribute to the diagram of Fig. 1b of the five-leg low-energy photon–neutrino processes (1), (2) and (3).

Concerning the second term in the Lagrangian (6), because of the antisymmetric nature of the coupling, the lepton-number violation is manifest: the family-type of the lepton running inside the loop is different from the one of the incoming neutrino. For example, if we assume that we are in a situation where $m_{\tilde{e}_L^k} \gg m_{\tilde{e}_R^k}$, so that the third term of the Lagrangian (6) is not relevant, then, when computing $\nu_e \gamma \to \nu_e \gamma\gamma$, the running lepton is either a muon or a tau and, of course, the muon contribution is the leading one. In this case, the most relevant contribution may come from the processes where the running lepton is an electron:

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$$\nu_\mu\gamma \rightarrow \nu_\mu\gamma\gamma \quad \text{or} \quad \nu_\tau\gamma \rightarrow \nu_\tau\gamma\gamma \ . \tag{7}$$

However, processes (7) are not relevant for the study of supernova dynamics since in this case, we are interested in the ν_e-type.

When we assume that $m_{\tilde{e}_L^k} \simeq m_{\tilde{e}_R^k}$, the third term in the Lagrangian (6) is always relevant since, for both processes (7) and the process $\nu_e\gamma \rightarrow \nu_e\gamma\gamma$, the running lepton inside the loop could be an electron. There is no constraint due to the antisymmetric nature of the coupling. It is the exchanged slepton that has to have a family-type different from the incoming neutrino.

Moreover, since the lepton number is violated in this model, we can have also the following flavour-changing neutrino transitions:

$$\nu_e\gamma \rightarrow \nu_\mu\gamma\gamma \ , \quad \nu_e\gamma \rightarrow \nu_\tau\gamma\gamma \quad \text{and} \quad \nu_\mu\gamma \rightarrow \nu_\tau\gamma\gamma \ . \tag{8}$$

Processes (8) could, in principle, play some role in the solar-neutrino puzzle, since the first two processes in (8) correspond to ν_e-type suppression. The numerical impact, though, could be insignificant on account of low cross sections.

Taking into account only the contributions where the running lepton inside the loop is an electron, we observe that it is straightforward to adapt the existing tools of the Standard Model computation [7], as long as we express the R_p-violating effective four-fermion operators (keeping only the charged slepton $\tilde{e}$ exchange) in the same $(V - A) \times (V - A)$ form [14]:

$$L_{\text{eff}} = \frac{\lambda_{ijk}^2}{2} \left[\left[\frac{1}{m_{\tilde{e}_R^k}^2} \left(\bar{\nu}_L^i \gamma^\mu \nu_L^i \right) \left(\bar{e}_L^j \gamma^\mu e_L^j \right) - \frac{1}{m_{\tilde{e}_L^i}^2} \left(\bar{\nu}_L^j \gamma^\mu \nu_L^j \right) \left(\bar{e}_R^k \gamma^\mu e_R^k \right) - \right.\right.$$
$$\left.\left. \frac{1}{m_{\tilde{e}_R^k}^2} \left(\bar{e}_L^i \gamma^\mu \nu_L^i \right) \left(\bar{\nu}_L^j \gamma^\mu e_L^j \right) + [i \leftrightarrow j] \right] \right] \ . \tag{9}$$

The contribution of the R_p processes in Fig. 1b thus has the same $(V - A) \times (V - A)$ structure as the W exchange in the Standard Model, so that the coupling g is replaced by the Yukawa coupling λ_{ijk} and the W-propagator by the slepton's $(M_W \rightarrow m_{\tilde{e}^k})$ in the computation of the amplitude. For simplicity, we will assume that sleptons are degenerate in mass, $m_{\tilde{e}^1} = m_{\tilde{e}^2} = m_{\tilde{e}^3}$.

To have an indication on the contribution of the processes involving the slepton exchange, we are assuming from now on a degenerate mass for $\tilde{e}_L$ and $\tilde{e}_R$.

In the following, we will make use of the notations of Ref. [7], where we have computed in the Standard Model the amplitudes and cross sections for the processes (1), (2) and (3),

the ν being of any family-type. Here, for the sake of comparison with the SM model results, we will concentrate on process (1), that is, $\nu_i \gamma \to \nu_j \gamma \gamma$. We will give in the following the different amplitudes according to the family-type of the neutrinos engaged in the process using the Lagrangian (6), in terms of the amplitudes computed in the Standard Model (M^{SM}) with the appropriate changes (the couplings, the exchanged-boson and the running lepton inside the loop). More precisely:

i) The coefficient[2] v_e corresponding to the Z exchange (Fig. 1a) is set to zero because, as already mentioned, only Fig. 1b will contribute here.

ii) The mass m_e in the SM amplitude is replaced by the one of the appropriate charged lepton running inside the loop.

iii) $1/\Delta_W \sim -1/M_W^2$ is replaced by $1/\Delta m_{\tilde{e}} \sim -1/m_{\tilde{e}}^2$, $m_{\tilde{e}}$ being degenerate.

We have then

$$M(\nu_e \gamma \to \nu_e \gamma \gamma) = \sum_{k=1,3} \lambda_{k11}^2 \frac{M^{\mathrm{SM}}(\nu_e \gamma \to \nu_e \gamma \gamma)}{g^2} [v_e \to 0, M_W \to m_{\tilde{e}}] \; , \tag{10}$$

$$M(\nu_\mu \gamma \to \nu_\mu \gamma \gamma) = \left(\sum_{k=1,3} \lambda_{21k}^2 + \sum_{i=1,3} \lambda_{i21}^2 \right) \frac{M^{\mathrm{SM}}(\nu_e \gamma \to \nu_e \gamma \gamma)}{g^2} [v_e \to 0, M_W \to m_{\tilde{e}}] \; , \tag{11}$$

$$M(\nu_\tau \gamma \to \nu_\tau \gamma \gamma) = \left(\sum_{k=1,3} \lambda_{31k}^2 + \sum_{i=1,3} \lambda_{i31}^2 \right) \frac{M^{\mathrm{SM}}(\nu_e \gamma \to \nu_e \gamma \gamma)}{g^2} [v_e \to 0, M_W \to m_{\tilde{e}}] \; . \tag{12}$$

The transitions given in (8) are also computed as follows:

$$M(\nu_e \gamma \to \nu_\mu \gamma \gamma) = \sum_{k=1,3} \lambda_{k11} \lambda_{k21} \frac{M^{\mathrm{SM}}(\nu_e \gamma \to \nu_e \gamma \gamma)}{g^2} [v_e \to 0, M_W \to m_{\tilde{e}}] \; , \tag{13}$$

$$M(\nu_e \gamma \to \nu_\tau \gamma \gamma) = \sum_{k=1,3} \lambda_{k11} \lambda_{k31} \frac{M^{\mathrm{SM}}(\nu_e \gamma \to \nu_e \gamma \gamma)}{g^2} [v_e \to 0, M_W \to m_{\tilde{e}}] \; , \tag{14}$$

$$M(\nu_\mu \gamma \to \nu_\tau \gamma \gamma) = \left(\sum_{k=1,3} \lambda_{21k} \lambda_{31k} + \sum_{k=1,3} \lambda_{k21} \lambda_{k31} \right) \frac{M^{\mathrm{SM}}(\nu_e \gamma \to \nu_e \gamma \gamma)}{g^2} [v_e \to 0, M_W \to m_{\tilde{e}}] \; . \tag{15}$$

In Eqs. (10–15), we are neglecting the contributions where the running lepton inside the loop is a muon or a tau. When the running lepton is an electron, in principle, since the

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[2]The coefficient v_e, which is directly related to the Z^0-exchange, is given by $v_e = -\frac{1}{2} + 2s_W^2$, where s_W (c_W) is the sine (cosine) of the Weinberg angle.

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computation is already done in the Standard Model and since the changes one needs to do in this case are global factors, the computation in the MSSM with $\not{R}_p$ is straightforward.

Using the definition

$$\frac{g^2}{8M_W^2} = \frac{G_F}{\sqrt{2}}$$

and taking the following couplings[3] as degenerate (see Eqs. (5)):

$$\lambda_{12k} = \lambda_{123} \sim 0.04, \quad \text{for} \quad k = 1, 2, 3 \ ,$$

$$\lambda_{131} = \lambda_{132} \sim 0.05, \quad \text{and} \quad \lambda_{133} \sim 0 \ ,$$

$$\lambda_{23k} = \lambda_{233} \sim 0.05, \quad \text{for} \quad k = 1, 2, 3 \ , \tag{16}$$

we find the following ratios of the cross section in the MSSM with $\not{R}_p$ versus the SM ones:

$$\frac{\sigma^{\mathrm{MSSM}}(\nu_e\gamma \to \nu_e\gamma\gamma)}{\sigma^{\mathrm{SM}}(\nu_e\gamma \to \nu_e\gamma\gamma)} \simeq 1 + \frac{\left(\sum\limits_{k=1,3} \lambda_{k11}^2\right)^2}{g^4} \frac{1}{(1+v_e)^2} \frac{M_W^4}{m_{\tilde{e}}^4} \sim 1 + 0.01\% \ , \tag{17}$$

$$\frac{\sigma^{\mathrm{MSSM}}(\nu_\mu\gamma \to \nu_\mu\gamma\gamma)}{\sigma^{\mathrm{SM}}(\nu_\mu\gamma \to \nu_\mu\gamma\gamma)} \simeq 1 + \frac{\left(\sum\limits_{k=1,3} \lambda_{21k}^2 + \sum\limits_{i=1,3} \lambda_{i21}^2\right)^2}{g^4} \frac{1}{v_e^2} \frac{M_W^4}{m_{\tilde{e}}^4} \sim 1 + 32\% \ , \tag{18}$$

$$\frac{\sigma^{\mathrm{MSSM}}(\nu_\tau\gamma \to \nu_\tau\gamma\gamma)}{\sigma^{\mathrm{SM}}(\nu_\tau\gamma \to \nu_\tau\gamma\gamma)} \simeq 1 + \frac{\left(\sum\limits_{k=1,3} \lambda_{31k}^2 + \sum\limits_{i=1,3} \lambda_{i31}^2\right)^2}{g^4} \frac{1}{v_e^2} \frac{M_W^4}{m_{\tilde{e}}^4} \sim 1 + 40\% \ . \tag{19}$$

When the final neutrino family-type is different from the initial one, we get:

$$\frac{\sigma^{\mathrm{MSSM}}(\nu_e\gamma \to \nu_\mu\gamma\gamma)}{\sigma^{\mathrm{SM}}(\nu_e\gamma \to \nu_e\gamma\gamma)} \simeq 1 + \frac{\left(\sum\limits_{k=1,3} \lambda_{k11}\lambda_{k21}\right)^2}{g^4} \frac{1}{(1+v_e)^2} \frac{M_W^4}{m_{\tilde{e}}^4} \sim 1 + 0.004\% \ , \tag{20}$$

$$\frac{\sigma^{\mathrm{MSSM}}(\nu_e\gamma \to \nu_\tau\gamma\gamma)}{\sigma^{\mathrm{SM}}(\nu_e\gamma \to \nu_e\gamma\gamma)} \simeq 1 + \frac{\left(\sum_{k=1,3} \lambda_{k11}\lambda_{k31}\right)^2}{g^4} \frac{1}{(1+v_e)^2} \frac{M_W^4}{m_{\tilde{e}}^4} \sim 1 + 0.003\% \ , \tag{21}$$

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[3]The bounds on the couplings λ_{ijk} have been rescaled according to the new degenerate slepton mass limit given by ALEPH [15].

$$\frac{\sigma^{\mathrm{MSSM}}(\nu_\mu\gamma \to \nu_\tau\gamma\gamma)}{\sigma^{\mathrm{SM}}(\nu_\mu\gamma \to \nu_\mu\gamma\gamma)} \simeq 1 + \frac{\left(\sum_{k=1,3}\lambda_{21k}\lambda_{31k} + \sum_{k=1,3}\lambda_{k21}\lambda_{k31}\right)^2}{g^4}\frac{1}{v_e^2}\frac{M_W^4}{m_{\tilde{e}}^4} \sim 1 + 0.03\% \, . \quad (22)$$

The limits on the slepton masses are taken from the latest ALEPH analysis [15], which gives $m_{\tilde{e}} \sim 80$ GeV for the degenerate mass of the sleptons; the bounds on the couplings λ_{ijk} are given by Eq. (16) [13]. The Standard Model cross sections are computed in Refs. [6,7].

One can see from Eqs. (20), (21) and (22) that an electron-neutrino transforms better in a muon-neutrino than in a tau-neutrino. However, in view of the smallness of the MSSM with $\not{R}_p$ cross sections, these transitions are unlikely to be relevant to the solar neutrino puzzle, although the neutrino energy is in the appropriate range (from 0.1 to ~ 15 MeV). On the other hand, in these processes, a muon-neutrino prefers to convert into a tau-neutrino rather than an electron-neutrino. This goes in the direction of the Superkamiokande result; however, our processes hold for energies much below the energy of the atmospheric neutrinos (> 1 GeV).

We have made the same computation in the case where $m_{\tilde{e}_L^k} \gg m_{\tilde{e}_R^k}$, so that the second term of the Lagrangian (6) is the most relevant and we have found that the $\not{R}_p$ MSSM cross sections are enhanced by a factor of the order of 10% at best, relative to the SM cross sections.

It is worth emphasizing that the same combinations of the couplings λ_{ijk} that contribute to flavour-changing neutrino transitions (8) might contribute as well to flavour-changing charged-lepton radiative transitions such as

$$\mu \to e\gamma \, , \quad \tau \to e\gamma \quad \text{and} \quad \tau \to \mu\gamma \, ,$$

decays that are very restricted by experiment [16]. In this model, a rough estimate of the cross section of, say, the $\mu \to e\gamma$ transition, which occurs through the sneutrino exchange, gives $\sim 3 \times 10^{-13}$ while the experimental bound is $\sim 10^{-11}$ [16].

Finally, the natural question that has to be addressed is how large the contribution of R-parity-conserving SUSY to the transitions $\nu\gamma \to \nu\gamma\gamma$ could be. The answer is that the cross sections then are smaller, since the particles in the loop are heavier, as is the case when the running fermion inside the loop is a chargino.

III. COMPUTATION OF $\gamma\nu \to \gamma\gamma\nu$ IN THE LEFT–RIGHT MODEL

Left–Right models are based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{\tilde{Y}}$ [12,17]. This is a natural framework to embed extra $W^\pm$ and Z gauge bosons that could be found in forthcoming colliders. For a more general case with more than one generation of $W'^\pm$ and Z', see Ref. [18]. We will keep the notation used in [17].

The leading contribution of this type of models to our processes consists mainly, of the substitution of the $W^\pm$ and Z propagators by the corresponding $W'^\pm$ and Z', and the modification of the standard couplings of $W^\pm$ and Z gauge bosons with fermions because of the mixing effect.

The charged-current gauge interactions of leptons are given by

$$\mathcal{L}_{CC} = \left(\begin{array}{cc} J_L^{\mu+} & J_R^{\mu+} \end{array} \right) \left(\begin{array}{c} W_{\mu L}^- \\ W_{\mu R}^- \end{array} \right) + \text{h.c.} = \left(\begin{array}{cc} J_W^{\mu+} & J_{W'}^{\mu+} \end{array} \right) \left(\begin{array}{c} W_\mu^- \\ W_\mu'^- \end{array} \right) + \text{h.c.} , \tag{23}$$

where

$$J_W^{\mu+} = \cos\alpha_\pm J_L^{\mu+} + \sin\alpha_\pm J_R^{\mu+} , \quad J_{W'}^{\mu+} = -\sin\alpha_\pm J_L^{\mu+} + \cos\alpha_\pm J_R^{\mu+} , \tag{24}$$

and the charged current associated with $SU(2)_L$ and $SU(2)_R$ lepton interactions are, respectively:

$$J_L^{\mu+} = \frac{g_L}{\sqrt{2}} \left(\overline{e_L}\gamma^\mu \nu_L \right) \ \text{and} \ J_R^{\mu+} = \frac{g_R}{\sqrt{2}} \left(\overline{e_R}\gamma^\mu \nu_R \right) \ . \tag{25}$$

From now on, we will call A and B the Standard Model amplitudes corresponding to Figs. 1a and 1b, respectively, which are given in [6].

From Eqs. (24) and (25), and using the definition of the mixing angle of Ref. [17], we find that type B diagrams (Fig. 1b), now called B^W, are modified through the mixing by:

$$B^W = \cos^2 \alpha_\pm B. \tag{26}$$

Also a new diagram (called $B^{W'}$), with a $W'^\pm$ instead of the $W^\pm$ boson, appears:

$$B^{W'} = \sin^2 \alpha_\pm \frac{M_W^2}{M_{W'}^2} B \ . \tag{27}$$

Similarly, the modification in the neutral-current gauge interactions (see [17]) gives rise to a modification of type A diagrams (Fig. 1a):

$$A^Z = \cos^2 \alpha_0 A \tag{28}$$

and to a new contribution, due to the Z', given by:

$$A^{Z'} = \sin^2 \alpha_0 \frac{M_Z^2}{M_{Z'}^2} A \ . \tag{29}$$

The total contribution of the equivalent set of diagrams of Eq. (15) in [6] or Eq. (10) in [7] is then

$$A_{123}^Z + A_{321}^Z + B_{123}^W + B_{321}^W + A_{123}^{Z'} + A_{321}^{Z'} + B_{123}^{W'} + B_{321}^{W'} = C_{LR}\Gamma_\mu L_1 \ , \tag{30}$$

where

$$C_{LR} = \frac{g_L^5 s_W^3}{2} \left[\left(\cos^2 \alpha_\pm + \sin^2 \alpha_\pm \frac{M_W^2}{M_{W'}^2} \right) + \rho \, v_e \left(\cos^2 \alpha_0 + \sin^2 \alpha_0 \frac{M_Z^2}{M_{Z'}^2} \right) \right] \frac{1}{M_W^2} \ , \tag{31}$$

where $\rho = M_W^2/c_W^2 M_Z^2$; L_1 and Γ_μ are given in [6]; L_1 can be evaluated in the large m_e limit, as in [6], or exactly as in [7]. The corresponding SM coefficient C_{SM} can be obtained trivially from Eq.(31) in the limit $\alpha_\pm \to 0$ and $\alpha_0 \to 0$.

Notice that the New Physics contribution enters, according to Eq. (30), as a multiplicative factor.

In order to evaluate C_{LR}, we will work in the small mixing-angle-approximation substituting $\cos \alpha_{\pm,0} \to 1$ and $\sin \alpha_{\pm,0} \to \alpha_{\pm,0}$. At this order, the ρ parameter that enters C_{LR} is given by [17]

$$\rho = 1 - \alpha_\pm^2 \left(\frac{M_{W'}^2 - M_W^2}{M_W^2} \right) + \alpha_0^2 \left(\frac{M_{Z'}^2 - M_Z^2}{M_Z^2} \right) \ . \tag{32}$$

Finally, we should define our input parameters. Concerning the sector of the model that affects our computation, we have only eight parameters, which we will choose to be α, M_W (or G_F), M_Z and m_e (as in the Standard Model) plus $x = g_R/g_L$, $M_{W'}$, $\alpha_\pm$ and α_0. In terms of these, in the small-mixing-angle approximation, we find, to the precision required for the evaluation of Eq. (31), that:

$$s_W^2 = 1 - \frac{M_W^2}{M_Z^2} - \frac{M_W^2}{M_Z^2} \left(\alpha_\pm^2 \left(\frac{M_{W'}^2 - M_W^2}{M_W^2} \right) - \alpha_0^2 \left(\frac{M_{Z'}^2 - M_Z^2}{M_Z^2} \right) \right) \tag{33}$$

and

$$M_{Z'}^2 = \frac{x^2 M_W^2 M_{W'}^2 - (M_Z^2 - M_W^2)^2}{x^2 M_W^2 - M_Z^2 + M_W^2} \tag{34}$$

and, of course, $g_L = \sqrt{4\pi\alpha}/s_W$. The stronger bounds [16] on our input parameters comes from the flavour-changing neutral-current FCNC (mainly the $K_L - K_S$ mass difference) [19,20], but they depend on the assumptions of the model (manifestly LR symmetric models [21], $g_L \neq g_R$ models [20], fermiophobic models [17,18], etc.). For instance, in manifestly symmetric Left-Right models with $g_L = g_R$, the bound on the mass of the W' is $M_{W'} \gtrsim 1.6$ TeV. The bounds on the mixing angles depend on the CP-violating phases of the theory; for small phases, it is $|\alpha_\pm| < 0.0025$; for large phases, it is $|\alpha_\pm| < 0.033$. If the constraint $g_L = g_R$ is relaxed, the bounds on $M_{W'}$ masses are much weaker. Finally, a fermiophobic model, which automatically guarantees the absence of FCNCs at tree level, allows for a relatively light $M_{W'}$ with no contradiction with experimental data.

For a typical set of values of the second half of input parameters ($x = g_R/g_L$, $M_{W'}$, $\alpha_\pm$ and α_0), taking into account the bounds on FCNCs for each model and the stringent bounds on ρ of Eq.(32), it is possible to obtain correction of at most a few per mille to the SM value of C_{SM}, which means also a few per mille enhancement in the cross sections.

IV. CONCLUSIONS

The cross sections computed in the $\not{R}_p$ MSSM are enhanced by a factor of the order of few 10% at best relative to the SM cross sections, while the correction in LR models is negligible.

Concerning the cosmological implications, in view of the small enhancements found, the conclusion that these low-energy five-leg processes are not relevant to the study of the neutrino decoupling temperature [7] remains unchanged.

The results found in these extensions of the SM will also enforce the conclusions made in [7] on the supernova question, that is: these low-energy processes should be taken into account in the supernova codes.

On the other hand, if the family-type of the neutrino changes in the transition, as it is allowed in the MSSM with $\not{R}_p$, the results found in this model then go in the direction of the actual conjecture, explaining results from both solar and atmospheric neutrino experiments, see Eqs. (20), (21) and (22). However, the cross sections are far too small to have any substantial effect on the neutrino fluxes.
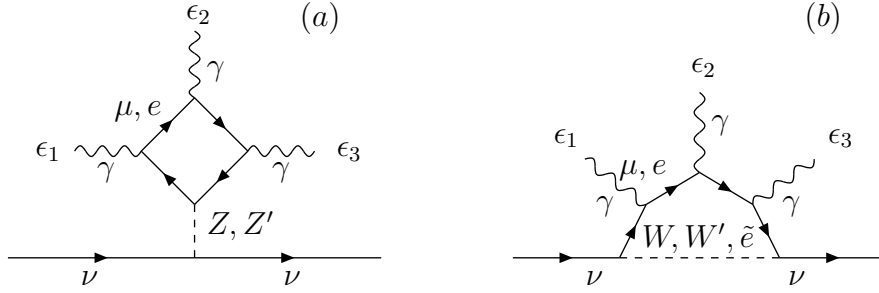
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Fig. 1: Leading diagrams to five-leg photon–neutrino low-energy processes. The couplings and the exchanged particles depend on the model.

REFERENCES

[1] H.Y. Chiu and P. Morrison, Phys. Rev. Lett. **5** (1960) 573; N. Van Hieu and E. P. Shabalin, Sov. Phys. JETP **17** (1963) 681; M.Harris, J. Wang and V.L. Teplitz, "Astrophysical effects of $\nu\gamma \to \nu\gamma\gamma$ and its crossed processes", astro-ph/9707113.

[2] J. Liu, Phys. Rev. **D44** (1991) 2879.

[3] D.A. Dicus and W.W. Repko, Phys. Rev. **D48** (1993) 5106.

[4] D.A. Dicus and W.W. Repko, Phys. Rev. Lett. **79** (1997) 569.

[5] A. Abada, J. Matias and R. Pittau, "Five-leg photon–neutrino interactions", hep-ph/9809418.

[6] A. Abada, J. Matias and R. Pittau, "Inelastic photon-neutrino interactions using an effective Lagrangian", hep-ph/9806383, Phys. Rev. **D59** (1999) 013008.

[7] A. Abada, J. Matias and R. Pittau, "Direct computation of inelastic photon–neutrino processes in the Standard Model", hep-ph/9808294, to appear in Nucl. Phys. **B**.

[8] C. N. Yang, Phys. Rev. **77** (1950) 242.

[9] H. Euler, Ann. Phys. **26** (1936) 398; W. Heisenberg and H. Euler, Z. Phys. **98** (1936) 714.

[10] D.A. Dicus, Chung Kao and W.W. Repko, " $\gamma\nu \to \gamma\gamma\nu$ and crossed processes at energies below m_W", hep-ph/9806499.

[11] For recent reviews on $\not{R}_p$ see for instance: G. Bhattacharyya, Nucl. Phys. Proc. Suppl. **52A** (1997)83 and "A brief review on R parity violating couplings", hep-ph/9709395 ; H. Dreiner, hep-ph/9707435, to be published in "Perspectives on Supersymmetry", ed. by G.L. Kane, World Scientific; R. Barbier et al., "Report of the group on the R-parity violation", hep-ph/9810232.

[12] J.C. Pati and A. Salam, Phys. Rev. **D10** (1974) 275; R.N. Mohapatra and J.C.Pati, Phys. Rev. **D11** (1975) 566 and 2558; G. Senjanovic and R.N. Mohapatra, Phys. Rev. **D12** (1975) 1502.

[13] See Ref. [14] and update in Ref. [11].

[14] V. Barger, G. F. Giudice and T. Han, Phys. Rev. **D 40** (1989) 2987.

13

[15] The ALEPH Collaboration (R. Barate et al.), Phys. Lett. **B433** (1998) 176.

[16] C. Caso et al., Eur Phys J **C3** (1998) 1.

[17] A. Donini, F. Feruglio, J. Matias and F. Zwirner, Nucl. Phys. **B507** (1997) 51.

[18] J. Matias and A. Vicini, hep-ph/9803278.

[19] G. Beall, M. Bander and A. Soni, Phys. Rev. Lett. **48** (1982) 848.

[20] P. Langacker and S. Uma Sankar, Phys. Rev. **D40** (1989) 1569; G. Barenboim, J. Bernabeu, J. Prades and M. Raidal, Phys. Rev. **D55** (1997) 4213.

[21] G. Senjanovic, Nucl. Phys. **B153** (1979) 334; M.A.B. Beg et al., Phys. Rev. Lett. **38** (1977) 1252.